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High Energy Quantum Teleportation Using Neutral Kaons

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We describe a scheme of stochastic implementations of quantum teleportation and entanglement swapping in terms of neutral kaons. In this scheme, the kaon whose state is to be teleported collides with one of the two entangled kaons in an Einstein-Podolsky-Rosen state. Subsequent detection of the outgoing particles of the collision completes the two-qubit projection on Alice side. There appear novel features, which connects quantum information science with fundamental laws of particle physics.

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Recent years witnessed the blossom of the subject of quantum information [1]. A major topic is quantum teleportation, the transmission of a quantum state without transporting the physical object through the intervening space [2], which has been implemented in optical, atomic and NMR systems [3]. Theoretical studies have also been made in a solid state electron-hole system [4] and in a non-inertial frame [5]. The idea of teleportation is even involved in studies on black hole evaporation [6]. In a closely related process called entanglement swapping, two qubits which never meet become entangled [7]. It has been implemented optically [8]. Virtually all areas of physics reigned by quantum mechanics have been explored for possibilities of implementing quantum information processes, but with an exception: high energy physics. Yet there have been many explorations on testing Bell theorem in terms of neutral kaons or B-mesons [9, 10, 11, 12, 14]. In this Letter, we make a theoretical proposal on high energy implementations of quantum teleportation and entanglement swapping by using neutral kaons. To our knowledge, it is the first proposal of quantum information processing in terms of massive elementary particles and in presence of particle decays. In this novel scheme, the kaon whose state is to be teleported collides with one of the entangled kaons, and the detection of the outgoing particles effectively realizes the two-qubit projection on Alice side, as required in a teleportation procedure. The projection basis is different from the Bell basis, but still contains the entangled state same as the original one shared between Alice and Bob. As a fundamental property of a massive elementary particle, the teleported degree of freedom, namely, being K^0 or $\bar{K}^0$, and its entanglement, are Lorentz invariant, in contrast with the case of spin. The cross-fertilization between quantum information science and high energy physics points to interesting new directions of research concerning some most fundamental aspects of information and matter, and may also lead to useful applications.

The neutral kaon K^0 is a meson composed of quarks d and $\bar{s}$, while its antiparticle $\bar{K}^0$ is composed of $\bar{d}$ and s [15]. Each of them is a pseudoscalar with $J^P = 0^-$, that is, the angular momentum is $J = 0$, while its intrinsic parity is negative, i.e. $P|K^0\rangle = -|K^0\rangle$ and $P|\bar{K}^0\rangle = -|\bar{K}^0\rangle$. They are also eigenstates of strangeness S with eigenvalues 1 and -1, respectively, as well as eigenstates of the component I_3 of isospin ($I = 1/2$), with eigenvalues $1/2$ and $-1/2$, respectively. They transform to each other by charge conjugation C, i.e. $C|K^0\rangle = \eta|\bar{K}^0\rangle$ and $C|\bar{K}^0\rangle = \eta'|K^0\rangle$, where η and η' are arbitrary phase factors, and are set to be -1 here. Under this convention,

$$|\bar{K}^0\rangle = CP|K^0\rangle,$$

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Hence the eigenstates of CP are

$$|K_1\rangle = \frac{1}{\sqrt{2}}(|K^0\rangle + |\bar{K}^0\rangle),$$

with $CP = 1$, and

$$|K_2\rangle = \frac{1}{\sqrt{2}}(|K^0\rangle - |\bar{K}^0\rangle),$$

with $CP = -1$. $|K^0\rangle$ and $|\bar{K}^0\rangle$, or $|K_1\rangle$ and $|K_2\rangle$, expand a two-dimensional Hilbert space. Another important basis is comprised of the mass eigenstates $|K_S\rangle$ and $|K_L\rangle$, with eigenvalues $\lambda_S = m_S - i\Gamma_S/2$ and $\lambda_L = m_L - i\Gamma_L/2$, where the subscripts "$S$" and "$L$" stand for the short and long life times of weak decays, with decay widths Γ_S and Γ_L, respectively. The mass difference is negligible, as $m_L - m_S = (3.483 \pm 0.006) \times 10^{-12} MeV$, while $m_L \approx m_S \approx m = 497.648 \pm 0.022 MeV$ [16], hence $m_L - m_S \approx 7.112 \times 10^{-15} m$. But the mean life times $1/\Gamma_S = (0.8953 \pm 0.0006) \times 10^{-10} s$ and $1/\Gamma_L =$

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$(5.114 \pm 0.021) \times 10^{-8}s$ differ significantly [16]. In terms of the proper time τ, the weak decay is described as

$$|K_S(\tau)\rangle = e^{-i\lambda_S\tau}|K_S\rangle,$$

$$|K_L(\tau)\rangle = e^{-i\lambda_L\tau}|K_L\rangle,$$

with $\hbar = c = 1$. $|K_L\rangle$ and $|K_S\rangle$ are related to CP and strangeness eigenstates as

$$\begin{aligned} |K_S\rangle &= \frac{1}{\sqrt{1+|\epsilon|^2}}(|K_1\rangle + \epsilon|K_2\rangle) \\ &= \frac{1}{\sqrt{|p|^2+|q|^2}}(p|K^0\rangle + q|\bar{K}^0\rangle), \\ |K_L\rangle &= \frac{1}{\sqrt{1+|\epsilon|^2}}(|K_2\rangle + \epsilon|K_1\rangle) \\ &= \frac{1}{\sqrt{|p|^2+|q|^2}}(p|K^0\rangle - q|\bar{K}^0\rangle), \end{aligned} \tag{1}$$

where ϵ is the very small parameter characterizing CP violation, and is of the order of 10^{-3}, $p = 1+\epsilon$, $q = 1-\epsilon$. In practice, in neglecting CP violation, one can set $p = q = 1$.

Note that the phase factors η and η' above can be chosen arbitrarily. If one adopts the convention $\eta = \eta' = 1$, then in Eq. (1), the expressions for $|K_S\rangle$ and $|K_L\rangle$ should be exchanged, and in the calculation results below, λ_S and λ_L should be exchanged. Anyway, $|K_S\rangle$ is always dominated by $CP = 1$ state, while $|K_L\rangle$ is always dominated by $CP = -1$ state.

It has been noted for a long time that from the strong decay of a vector meson ϕ or from the annihilation of a proton-antiproton pair, a $K^0\bar{K}^0$ pair can be created in an entangled Einstein-Podolsky-Rosen (EPR) state [10]

$$|\Psi_-\rangle \equiv \frac{1}{\sqrt{2}}(|K^0\rangle|\bar{K}^0\rangle - |\bar{K}^0\rangle|K^0\rangle).$$

ϕ mesons can be generated in electron-positron annihilation with center of mass energy about $1GeV$, as done in ϕ factories. Similar EPR state can be produced in $B^0\bar{B}^0$ pair from $\Upsilon(4S)$ resonance, which can be generated in electron-positron annihilation at about $10GeV$, as in B factories. A lot of discussions were made on how to employ the entangled $K^0\bar{K}^0$ or $B^0\bar{B}^0$ pair to test nonlocality or Bell theorem [9, 10, 11]. Experimentally, EPR correlation has been confirmed in $K^0\bar{K}^0$ pairs produced in proton-antiproton annihilation in the CPLEAR detector in CERN [12], in $K^0\bar{K}^0$ pairs produced in ϕ decay in the KLOE detector in DAΦNE [13], as well as in $B^0\bar{B}^0$ pairs produced in the Belle detector in the KEKB electron-position collider [14].

Our proposal is to let Alice and Bob share an entangled $K^0\bar{K}^0$ (or $B^0\bar{B}^0$) pair, denoted as a and b, and use it to teleport to b an unknown state of another kaon (or B-meson) c, be it in a pure state or entangled with another system. Our scenario involves collision between kaons c and a, and subsequent measurement of the outgoing particles of $c - a$ collision.

At laboratory time $t = t_y = 0$, an entangled $K^0\bar{K}^0$ pair a and b is created as $|\Psi_-\rangle$. Thus

$$\begin{aligned} |\Psi_{ab}(0)\rangle &= \frac{1}{\sqrt{2}}(|K^0\rangle_a|\bar{K}^0\rangle_b - |\bar{K}^0\rangle_a|K^0\rangle_b) \\ &= \frac{r}{\sqrt{2}}(|K_L\rangle_a|K_S\rangle_b - |K_S\rangle_a|K_L\rangle_b), \end{aligned}$$

where $r = (|p|^2 + |q|^2)/2pq \approx 1$. Up to r, the singlet in strangeness basis is also a singlet in mass basis even though CP violation is taken into account. For convenience, the observers, or Alice and Bob, which are particle detectors, stay in the laboratory frame, which is supposed to coincide with the center of mass frame of a and b; the generalization to otherwise case is straightforward.

After creation, the kaons naturally decay under weak interaction. It can be found that

$$|\Psi_{ab}(t)\rangle = M(t)|\Psi_-\rangle_{ab},$$

where $M(t) = \exp[-i(\lambda_S + \lambda_L)\gamma_b t]$. γ_i is the Lorentz factor $1/\sqrt{1 - v_i^2}$ for particle i with velocity v_i. It has been assumed that $\gamma_a = \gamma_b$. It can be estimated that with such decay widths and with the center of mass energy of about $1GeV$, the process should be completed within a few centimeters from the source in the laboratory frame, as indeed so in the CERN experiment [12]. Interestingly, $|\Psi_{ab}(t)\rangle$ is Lorentz invariant, as the kaons are spinless pseudoscalars. The Lorentz invariance of this entanglement is an advantage over the spin entanglement, which is not Lorentz invariant in general [19].

Next we consider the third kaon c generated at time t_z as

$$|\Psi_c(t_z)\rangle = \alpha|K^0\rangle_c + \beta|\bar{K}^0\rangle_c,$$

which may be unknown. For $t \geq t_z$,

$$|\Psi_c(t)\rangle = F(t)|K^0\rangle_c + G(t)|\bar{K}^0\rangle_c, \tag{2}$$

with $F(t) = [(\alpha + \beta p/q)e^{-i\lambda_S\gamma_c(t-t_z)} + (\alpha - \beta p/q)e^{-i\lambda_L\gamma_c(t-t_z)}]/2$, $G(t) = [(\alpha q/p + \beta)e^{-i\lambda_S\gamma_c(t-t_z)} - (\alpha q/p - \beta)e^{-i\lambda_L\gamma_c(t-t_z)}]/2$.

The state of the three particles is thus

$$|\Psi_{cab}(t)\rangle = |\Psi_c(t)\rangle \otimes |\Psi_{ab}(t)\rangle.$$

Consider the following basis states of $c - a$, which are eigenstates of P, S and I: $|\phi_1\rangle_{ca} \equiv |K^0K^0\rangle$ with $P = 1$, $S = 2$, $I = 1$; $|\phi_2\rangle_{ca} \equiv |\bar{K}^0\bar{K}^0\rangle$ with $P = 1$, $S = -2$, $I = 1$; $|\phi_3\rangle_{ca} \equiv |\Psi_+\rangle \equiv \frac{1}{\sqrt{2}}(|K^0\rangle|\bar{K}^0\rangle + |\bar{K}^0\rangle_c|K^0\rangle)$, with $P = 1$, $S = 0$, $I = 1$; and $|\phi_4\rangle_{ca} \equiv |\Psi_-\rangle$ with $P = -1$, $S = 0$, $I = 0$. For the reason which will be clear shortly, we rewrite $|\Psi_{cab}(t)\rangle$ in terms of these four basis states of $c - a$ as

$$\begin{aligned} |\Psi_{cab}(t)\rangle = \frac{M(t)}{2}\{ &\sqrt{2}F(t)|\phi_1\rangle_{ca}|\bar{K}^0\rangle_b \\ &-\sqrt{2}G(t)|\phi_2\rangle_{ca}|K^0\rangle_b \\ &-|\phi_3\rangle_{ca}[F(t)|K^0\rangle_b - G(t)|\bar{K}^0\rangle_b] \\ &-|\phi_4\rangle_{ca}[F(t)|K^0\rangle_b + G(t)|\bar{K}^0\rangle_b]\}. \end{aligned} \tag{3}$$

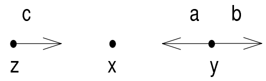

FIG. 1: Scheme of quantum teleportation using an entangled kaon pair a and b generated from a source y. Another kaon c comes from a source z. a and c fly collinearly and towards each other, thus collide at a certain time t_x at a position x.

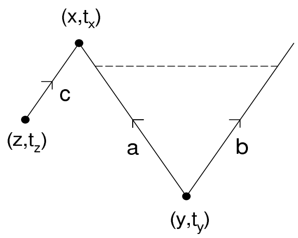

FIG. 2: Spacetime diagram of the kaon teleportation. The horizontal direction represents the position while the upward direction represents time flow. The broken line represents the entanglement. (z,t_z), (y,t_y) and (x,t_x) are spacetime coordinates of the generation of c, the generation of $a-b$ entangled pair, and the $c-a$ collision, respectively, in the laboratory frame.

We design the set-up in such a way that a and c fly in opposite directions and towards each other, hence they destine to collide at a certain position x at a certain time t_x (FIG. 1 shows the scheme, FIG. 2 is the spacetime diagram).

Upon collision, c and a become an interacting whole. The effect of collision can be represented as a unitary transformation $\mathcal{S}$ on $c-a$. The brief and negligible time duration δ of the collision is much shorter than the life times of weak decay, thus we ignore the decay of kaon b during the negligible interval of $c-a$ collision. Therefore, through $c-a$ collision, the state of the three kaons becomes

$$\begin{aligned}|\Psi_{cab}(t_x+\delta)\rangle \;=\; & \tfrac{M(t_x)}{2}\{\sqrt{2}F(t_x)\mathcal{S}|\phi_1\rangle_{ca}|\bar{K}^0\rangle_b \\ & -\sqrt{2}G(t_x)\mathcal{S}|\phi_2\rangle_{ca}|K^0\rangle_b \\ & -\mathcal{S}|\phi_3\rangle_{ca}[F(t_x)|K^0\rangle_b - G(t_x)|\bar{K}^0\rangle_b] \\ & -\mathcal{S}|\phi_4\rangle_{ca}[F(t_x)|K^0\rangle_b + G(t_x)|\bar{K}^0\rangle_b]\}.\end{aligned} \tag{4}$$

$\{\mathcal{S}|\phi_i\rangle_{ca}\}$ is also a basis of $c-a$ system. Furthermore, as $c-a$ collision is governed by strong interaction, $\mathcal{S}$ conserves S, P and I, thus $\mathcal{S}|\phi_i\rangle_{ca}$ $(i=1,2,3,4)$ is still an eigenstate of S, P and I, with the same eigenvalues as those for $|\phi_i\rangle_{ca}$.

The two-particle projection of $c-a$ is completed when the outgoing particles of $c-a$ collision are detected. The particle detector, close and around the collision point, playing the role of Alice, detects particles with specific values of S, P and I, thus will have projected $c-a$ to one of the four eigenstates $\mathcal{S}|\phi_i\rangle_{ca}$. Note that it may be particles other than kaons that are detected in $\mathcal{S}|\phi_i\rangle_{ca}$. But this does not matter, as S, P and I are conserved by $\mathcal{S}$. The detector may be based on strong interaction with absorbers, and constructed in a way similar to that in CPLEAR [12].

According to standard quantum theory, upon measurement (detection), the state instantaneously projects to an eigenstate of the observable [17]. The reference frame can be chosen arbitrarily, but once it is chosen, the instantaneous projection needs to be consistently made [18].

The probability of projection to $\mathcal{S}|\phi_i\rangle_{ca}$ is calculated as $\langle\Psi_{cab}(t_x+\delta)\mathcal{S}|\phi_i\rangle\langle\phi_i|\mathcal{S}^\dagger\Psi_{cab}(t_x+\delta)\rangle$. It is calculated that for $i=1,2,3,4$, the probabilities are $|M(t_x)|^2|F(t_x)|^2/2$, $|M(t_x)|^2|G(t_x)|^2/2$, $|M(t_x)|^2[|F(t_x)|^2+|G(t_x)|^2]/4$, and $|M(t_x)|^2[|F(t_x)|^2+|G(t_x)|^2]/4$, respectively.

Thus the two-qubit operation by Alice in the teleportation protocol is effectively realized. Conditioned on P, S and I of the detected outgoing particles of $c-a$ collision, the state of b is known to be correspondingly in one of the four states $|\bar{K}^0\rangle_b$, $|K^0\rangle_b$, $[F(t_x)|K^0\rangle_b - G(t_x)|\bar{K}^0\rangle_b]/\sqrt{|F(t_x)|^2+|G(t_x)|^2}$, $[F(t_x)|K^0\rangle_b+G(t_x)|\bar{K}^0\rangle_b]/\sqrt{|F(t_x)|^2+|G(t_x)|^2}$. A noteworthy point is that despite the decay, the state after projection should still be normalized; the decay effect has been taken into account in the projection probability. As P and S are already sufficient to distinguish $\mathcal{S}|\phi_i\rangle_{ca}$ for different i's, it is not necessary to consider I.

It is difficult to implement subsequent precise one-bit unitary transformations on b particle, which is a part of the conventional scheme of teleportation. Hence we suggest to adopt a stochastic strategy, as follows. With negligible time delay, upon receiving the communication of the projection result of Alice, Bob decides whether to retain or abandon b particle according to whether the state is what he needs. In actual experiments, the classical communication and the subsequent conditional operation can be realized by an automatic control system. As in the usual scenario of teleportation, suppose Bob wishes to obtain $[F(t_x)|K^0\rangle_b + G(t_x)|\bar{K}^0\rangle_b]/\sqrt{|F(t_x)|^2+|G(t_x)|^2}$. According to the expansion in Eq. (4), Bob should retain the particle if and only if the projection result of c and a is $\mathcal{S}|\Psi_-\rangle_{ca}$.

The four possible projection results at $t_x+\delta$ lead to different values of strangeness ratio $\xi(t\geq t_x+\delta)$ of b particle, which can be experimentally verified in terms of reaction with nuclear matter. For $|\Psi(t\geq t_x+\delta)\rangle_b = f(t)|K^0\rangle_b + g(t)|\bar{K}^0\rangle_b$, $\xi(t)\equiv|f(t)|^2/g(t)|^2$. Many runs of the procedure are needed to measure this quantity. If irrespective of the projection results of $c-a$, b particles in the different runs of the experiment are all considered in measuring $\xi(t)$, then $\xi(t)$ should be calculated by using $|\Psi_{cab}(t)\rangle$, consequently $\xi(t)=1$. In contrast, if only b particles in those runs of the experiment with a certain projection result of $c-a$ are considered in measuring $\xi(t)$, then $\xi(t)$ is calculated by using the corresponding projected state of b. Denote the state of b following the projection as $\alpha(t_x)|K^0\rangle+\beta(t_x)|\bar{K}^0\rangle$. Its subsequent evolution is then similar to Eq. (2), with t_z substituted by

$t_x+\delta$, γ_c by γ_b, α by $\alpha(t_x)$, β by $\beta(t_x)$. It can be found that $\xi(t) = |\alpha(t_x)(e^{-\Gamma_S\tau/2}+e^{-\Gamma_L\tau/2})+\beta(t_x)(e^{-\Gamma_S\tau/2}-e^{-\Gamma_L\tau/2})|^2/|\alpha(t_x)(e^{-\Gamma_S\tau/2}-e^{-\Gamma_L\tau/2})+\beta(t_x)(e^{-\Gamma_S\tau/2}+e^{-\Gamma_L\tau/2})|^2$, where $\tau=\gamma_b(t-t_x-\delta)$. For each of the four projection cases, $\xi(t)$ is very different from 1. For example, if the teleportation is successful, i.e. the projection result of $c-a$ is $|\Psi_-\rangle_{ca}$, then $\xi(t)$ is given by substituting $\alpha(t_x) = F(t_x)/\sqrt{|F(t_x)|^2+|G(t_x)|^2}$, $\beta(t_x) = G(t_x)/\sqrt{|F(t_x)|^2+|G(t_x)|^2}$. Like delay choice in entanglement swapping [20], the projection results of $c-a$ can even be revealed only after all the experiments are finished, and are then used to sort the runs of the procedure to four subensembles corresponding to the four projection results.

Now we consider a similar stochastic implementation of entanglement swapping. In addition to $|\Psi_-\rangle_{ab}$ generated at time $t_y=0$, another kaon pair d and c is generated as $|\Psi_-\rangle_{dc}$ at time t_z. Similar to $|\Psi_{ab}(t)\rangle$, we have

$$|\Psi_{dc}(t)\rangle = M'(t-t_z)|\Psi_-\rangle_{dc},$$

where $M'(t-t_z)=\exp[-i(\lambda_S+\lambda_L)\gamma_d(t-t_z)]$, supposing $\gamma_c=\gamma_d$. Thus the state of the four particles is

$$|\Psi_{dcab}(t)\rangle = M'(t-t_z)M(t)|\Psi_-\rangle_{dc}|\Psi_-\rangle_{ab}.$$

In terms of the $c-a$ eigenstates of P, S and I, as given above, $|\Psi_{dcab}(t)\rangle$ can be written as

$$\begin{aligned}|\Psi_{dcab}(t)\rangle \;=\; & \tfrac{M'(t-t_z)M(t)}{2}(|\Psi_+\rangle_{ca}|\Psi_+\rangle_{db}-|\Psi_-\rangle_{ca}|\Psi_-\rangle_{db}\\ & -|K^0K^0\rangle_{ca}|\bar{K}^0\bar{K}^0\rangle_{db}-|\bar{K}^0\bar{K}^0\rangle_{ca}|K^0K^0\rangle_{db}).\end{aligned}$$

Similar to the above scenario of pure state teleportation, we let c and a fly towards each other to collide at a certain time t_x (FIG. 3 and FIG. 4). Within a negligible time interval δ, the collision, effecting a unitary transformation $\mathcal{S}$ on $c-a$, evolves $|\Psi_{dcab}(t_x)\rangle$ to

$$\begin{aligned}|\Psi_{dcab}(t_x+\delta)\rangle \;=\; & \tfrac{M'(t_x-t_z)M(t_x)}{2}(\mathcal{S}|\Psi_+\rangle_{ca}|\Psi_+\rangle_{db}\\ & -\mathcal{S}|\Psi_-\rangle_{ca}|\Psi_-\rangle_{db}\\ & -\mathcal{S}|K^0K^0\rangle_{ca}|\bar{K}^0\bar{K}^0\rangle_{db}\\ & -\mathcal{S}|\bar{K}^0\bar{K}^0\rangle_{ca}|K^0K^0\rangle_{db}).\end{aligned}$$

Then, in detecting outgoing particles from $c-a$ collision, c and a are projected to one of the four states $\mathcal{S}|\Psi_+\rangle_{ca}$, $\mathcal{S}|\Psi_-\rangle_{ca}$, $\mathcal{S}|K^0K^0\rangle_{ca}$ and $\mathcal{S}|\bar{K}^0\bar{K}^0\rangle_{ca}$, and then P, S and I are measured. Correspondingly d and b are projected to $|\Psi_+\rangle_{ca}$, $|\Psi_-\rangle_{ca}$, $|K^0K^0\rangle_{ca}$ and $|\bar{K}^0\bar{K}^0\rangle_{ca}$, respectively, each with probability $|M'(t_x-t_z)M(t_x)|^2/4$. Again, the projection result is revealed by P, S and I of the outcomes of $c-a$ collision, according to which Bob chooses to retain or abandon b particle.

The effect of entanglement swapping can be verified by measuring the strangeness asymmetry between b and d, defined as $A(t) = [p_{diff}(t)-p_{same}(t)]/[p_{diff}(t)+p_{same}(t)]$, where $p_{diff}(t)$ and $p_{same}(t)$ are, respectively, the probabilities

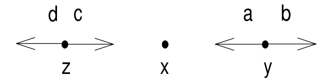

FIG. 3: Scheme of entanglement swapping using entangled kaon pair $a-b$ generated from a source y, and another entangled kaon pair $c-d$ generated from a source x, both in EPR states. a and c fly collinearly and towards each other, thus collide at a certain time t_x at a position x.

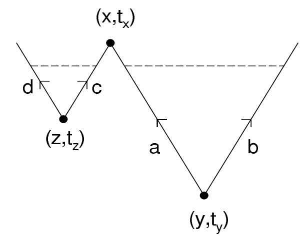

FIG. 4: Spacetime diagram of the entanglement swapping of kaons.

for b and d to have different and same strangeness values [12]. Many runs of the experiment are needed to experimentally determine $A(t)$. If all the $d-b$ pairs in different runs are considered, irrespective of the projection results of $c-a$, then $A(t)=0$, as calculated from $|\Psi_{dcab}(t)\rangle$. In contrast, if only the $d-b$ pairs corresponding to a certain projection result of $c-a$ collision are considered, then $A(t)$ is calculated by using the corresponding projected state of d and b. For example, if at $t=t_x+\delta$, c and a are projected to $\mathcal{S}|\Psi_-\rangle_{ca}$, i.e. the entanglement swapping is successful, then for $t\geq t_x+\delta$, $|\Psi(t\geq t_x+\delta)\rangle_{db} = |\Psi(t\geq t_x+\delta)\rangle_{db} = g_1|K^0\rangle_d|K^0\rangle_b + g_2|K^0\rangle_d|\bar{K}^0\rangle_b + g_3|\bar{K}^0\rangle_d|K^0\rangle_b + g_4|\bar{K}^0\rangle_d|\bar{K}^0\rangle_b$, where $g_1 = (e^{-i(\lambda_L\tau_d+\lambda_S\tau_b)} - e^{-i(\lambda_S\tau_d+\lambda_L\tau_b)})p/2\sqrt{2}q$, $g_4 = -(e^{-i(\lambda_L\tau_d+\lambda_S\tau_b)} - e^{-i(\lambda_S\tau_d+\lambda_L\tau_b)})q/2\sqrt{2}p$, $g_2 = (e^{-i(\lambda_L\tau_d+\lambda_S\tau_b)} + e^{-i(\lambda_S\tau_d+\lambda_L\tau_b)})/2\sqrt{2}$, $g_3=-g_2$, where $\tau_d=\gamma_d(t-t_x-\delta)$, $\tau_b=\gamma_b(t-t_x-\delta)$. Consequently $A(t\geq t_x+\delta)\approx 2e^{-(\Gamma_S+\Gamma_L)(\tau_d+\tau_b)/2}/(e^{-(\Gamma_L\tau_d+\Gamma_S\tau_b)}+e^{-(\Gamma_L\tau_d+\Gamma_S\tau_b)})$. If at $t=t_x+\delta$, c and a are projected to $\mathcal{S}|\Psi_+\rangle_{ca}$, then $|\Psi(t\geq t_x+\delta)\rangle_{db} = g_1|K^0\rangle_d|K^0\rangle_b + g_2|K^0\rangle_d|\bar{K}^0\rangle_b + g_3|\bar{K}^0\rangle_d|K^0\rangle_b + g_4|\bar{K}^0\rangle_d|\bar{K}^0\rangle_b$, where $g_1 = (e^{-i\lambda_S(\tau_d+\tau_b)} - e^{-i\lambda_L(\tau_d+\tau_b)})p/2\sqrt{2}q$, $g_4 = (e^{-i\lambda_S(\tau_d+\tau_b)} - e^{-i\lambda_L(\tau_d+\tau_b)})q/2\sqrt{2}p$, $g_2 = g_3 = (e^{-i\lambda_S(\tau_d+\tau_b)}+e^{-i\lambda_L(\tau_d+\tau_b)})/2\sqrt{2}$. Hence $A(t\geq t_x+\delta)\approx 2e^{-(\Gamma_S+\Gamma_L)(\tau_d+\tau_b)/2}/(e^{-\Gamma_S(\tau_d+\tau_b)}+e^{-\Gamma_L(\tau_d+\tau_b)})$. It can also be calculated that for $c-a$ projection to $\mathcal{S}|K^0K^0\rangle_{ca}$ or $\mathcal{S}|\bar{K}^0\bar{K}^0\rangle_{ca}$ at $t_x+\delta$, $A(t\geq t_x+\delta)\approx -4e^{-(\Gamma_S+\Gamma_L)\tau_b}/(e^{-2\Gamma_S\tau_b}+2e^{-(\Gamma_S+\Gamma_L)\tau_b}+e^{-2\Gamma_L\tau_b})$ in case $\gamma_b=\gamma_d$; the expressions for the case $\gamma_b\neq\gamma_d$ are too cumbersome to be included here. Again, the projection results of $c-a$ can even be revealed only after all the runs of the experiment are finished, and are then used to sort the runs to four subensembles corresponding to the four projection results.

As noted in the original paper on teleportation [2], tele-

portation can be made on a qubit c arbitrarily entangled with any other system d, with both the usual teleportation of a pure state and the entanglement swapping being special cases. The detail is the following. Suppose an arbitrary unknown state $|\Psi_0\rangle$ is shared between the qubit c and another system d. One can always write $|\Psi_0\rangle_{dc} = \gamma|w_1\rangle_d \uparrow_c + \rho|w_2\rangle_d \downarrow_c$, where $|w_1\rangle_d$ and $|w_2\rangle_d$ are normalized, but not necessarily orthogonal. Suppose a and b share a Bell state, say $|\Psi_-\rangle_{ab}$. Then $|\Psi\rangle_{dcab}$ can be written as

$$\begin{aligned}|\Psi\rangle_{dcab} \;=\; & \tfrac{1}{2}|\Phi_+\rangle_{ca}(\gamma|w_1\rangle_d \downarrow_b - \rho|w_2\rangle_d \uparrow_b) \\ & +\tfrac{1}{2}|\Phi_-\rangle_{ca}(\gamma|w_1\rangle_d \downarrow_b + \rho|w_2\rangle_d \uparrow_b) \\ & -\tfrac{1}{2}|\Psi_+\rangle_{ca}(\gamma|w_1\rangle_d \uparrow_b - \rho|w_2\rangle_d \downarrow_b) \\ & -\tfrac{1}{2}|\Psi_-\rangle_{ca}(\gamma|w_1\rangle_d \uparrow_b + \rho|w_2\rangle_d \downarrow_b).\end{aligned}$$

A Bell measurement is performed on c and a, with the resulting state $|\Phi_+\rangle_{ca}$, $|\Phi_-\rangle_{ca}$, $|\Psi_+\rangle_{ca}$ or $|\Psi_-\rangle_{ca}$. Correspondingly, the state of $d-b$ projects to $-i(\sigma_y)_b|\Psi_0\rangle_{db}$, $(\sigma_x)_b|\Psi_0\rangle_{db}$, $(\sigma_z)_b|\Psi_0\rangle_{db}$ and $|\Psi_0\rangle_{db}$, respectively, where $(\sigma_i)_b$ represents Pauli operation acting on b. Therefore, depending on the measurement result of $e-a$, Bob can correspondingly perform on b operation σ_y or σ_x or σ_z or make no operation. $|\Psi_0\rangle$ is then teleported to the same state shared between d and b.

Such a general teleportation can also be implemented in high energy mesons, in a way similar to the above scheme. To prepare an arbitrarily entangled kaon pair, one can, for example, let one or both particles of the kaon pair $c-d$, generated as the EPR state $|\Psi_-\rangle_{dc}$, pass through regeneration materials, which change the superposition coefficients [15]. Thus one obtains an unknown state, which is supposed to decay to $|\Psi_0(t_x)\rangle_{dc} = \gamma|w_1\rangle_d|K^0\rangle_c + \rho|w_2\rangle_d|\bar{K}^0\rangle_c$ at time t_x when c and a collide. With kaon pair $a-b$ prepared at $t=0$ in $|\Psi_-\rangle_{ab}$, one can write

$$\begin{aligned}|\Psi_{dcab}(t_x)\rangle \;=\; & \tfrac{M(t_x)}{2}[\sqrt{2}\gamma|K^0K^0\rangle_{ca}|w_1\rangle_d|\bar{K}^0\rangle_b \\ & -\sqrt{2}\rho|\bar{K}^0\bar{K}^0\rangle_{ca}|w_2\rangle_d|K^0\rangle_b \\ & -|\Psi_+\rangle_{ca}(|w_1\rangle_d|K^0\rangle_b - \rho|w_2\rangle_d|\bar{K}^0\rangle_b) \\ & -|\Psi_-\rangle_{ca}(\gamma|w_1\rangle_d|K^0\rangle_b + \rho|w_2\rangle_d|\bar{K}^0\rangle_b)].\end{aligned}$$

Afterwards, one proceeds in the same way as the above scheme of entanglement swapping. One can choose to retain the $d-b$ pair only if the detected outgoing particles of $c-a$ collision are with $P=-1$, $S=0$ and $I=0$, i.e. $c-a$ is projected to $\mathcal{S}|\Psi_-\rangle_{ca}$. In this way, one obtains $\gamma|w_1\rangle_d|K^0\rangle_b + \rho|w_2\rangle_d|\bar{K}^0\rangle_b$.

We see no particular obstacles in actually implementing our proposal. Especially, with the previous experimental experiences in nonlocality study [12, 13, 14], it looks feasible to implement teleportation in these places. Compared with the experiments on detecting CP violation and on testing Bell theorem, we only need the additional facilities of realizing kaon collision and the subsequent detection of the outgoing particles. We do not need either the precision as high as that in detecting CP violation or the particular arrangements required by the subtle argument of Bell theorem. Hence in some aspects, implementing the teleportation scheme here may be easier than detecting CP violation and testing Bell theorem. It seems easier to implement entanglement swapping than teleportation of either a pure kaon state or a kaon entangled with another kaon in an unknown state, as only two pairs of kaons in $|\Psi_-\rangle$ need to be prepared for entanglement swapping.

It is interesting to study high energy processes involving a single copy of particles as employed in quantum information protocols, such that the quantum nature of the processes can be more manifested. On the other hand, present high energy experiments often employ a beam of particles consisting of a group of particles prepared in the same state. Hence many runs of the same procedure mentioned in the above discussions can actually be done altogether simultaneously, as in the CPLEAR experiment [12]. In detecting $c-a$ collision and in analyzing the correlation between b and d particles, one needs to find the correspondence between events of the entangled particles in a same copy of state. This can be achieved by analyzing the particle trajectories and momenta, as done in CPLEAR experiment. The details of the collision between the kaons and the subsequent detection need further studies. Moreover, high energy processes in presence of entanglement with distant particles, as depicted in our spacetime diagrams, pose a new subject worth detailed investigations.

To summarize, we have described a scheme of stochastic implementations of quantum teleportation and entanglement swapping using neutral kaons. This work connects quantum information science to particle physics. The neutral kaon whose state is to be teleported collides with a neutral kaon which is entangled with another one in an EPR state. The detection of the outgoing particles of the collision completes the projection on Alice's side to an eigenstate of P, S and I, which is conserved in the collision, as governed by strong interaction. Conditioned on this projection, teleportation or entanglement swapping can be made stochastically. We also envisage verification schemes based on strangeness measurements. We expect our discussion to open up researches on high energy quantum information, which can stimulate particle physics to study processes involving entanglement, projection and decoherence, like similar developments in other areas of physics. They also furnish a fruitful playground for the extension of the notions of quantum information science to regimes of relativity and high energy [5, 6, 18, 19, 21]. There are some profound implications. For instance, in conventional teleportation, it is only the state, rather than the particle which carries the state, that is teleported. In high energy physics, the particle itself, e.g. K^0 or $\bar{K}^0$, represents a state of the quantum field, hence can be teleported, as demonstrated here in the example of neutral kaons.

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